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Quantum Restoration of the $\mathrm{U}(1)_Y$ Symmetry in Dynamically Broken SUSY-GUT's

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Abstract

We propose a supersymmetric hypercolor $\mathrm{SU}(3)_H$ gauge theory interacting strongly at the grand unification scale, in which the hyperquark condensation breaks $\mathrm{SU}(5)_{GUT}$ down to $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L$ without unbroken $\mathrm{U}(1)_Y$ at the classical level. However, we show that the broken $\mathrm{U}(1)_Y$ symmetry is restored by quantum mechanical effects and hence there remains the standard-model gauge symmetry at the electroweak scale. The dynamics of the strong interactions also produces naturally a pair of massless Higgs doublets. In addition to these Higgs doublets, we have a pair of massless singlets which contributes to the renormalization-group equations of gauge coupling constants and hence affects the GUT unification. We discuss a simple solution to this problem.

1. Introduction

An $\mathrm{SU}(2)_L$-doublet Higgs boson of mass of the order of hundred GeV is an inevitable ingredient in the standard electroweak gauge model. In the grand unified theories (GUT's) [1] one must require an extremely precise fine-tuning among various parameters in order to obtain the light doublet Higgs boson. The situation is not improved much even if one assumes supersymmetry (SUSY) [2]. Therefore, the necessity of the fine-tuning of parameters is regarded as a crucial drawback in the GUT scenario, although the recent high-precision measurements on the gauge coupling constants have strongly supported the SUSY extension of the GUT's [3].

In recent papers [4, 5, 6] we have proposed a SUSY gauge theory based on $\mathrm{SU}(3)_H \times \mathrm{U}(1)_H$ with six flavors of quarks which interact strongly at the GUT scale and shown that the dynamics of this theory produces a pair of massless Higgs doublets. Thus, we have a natural explanation for the presence of light Higgs doublets at the electroweak scale in the SUSY-GUT's.

However, we have also found that the hyperquark condensation breaks the GUT gauge group $\mathrm{SU}(5)_{GUT}$ down to $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L$. Therefore, we need to introduce an extra $\mathrm{U}(1)_H$ to have the standard-model gauge group unbroken at the low energies [4, 5, 6]. The presence of $\mathrm{U}(1)_H$ above the GUT scale is nothing wrong phenomenologically, but it would be more interesting if we did not need to introduce the extra $\mathrm{U}(1)_H$.

In this paper we show that introduction of a singlet field Φ into the previous model solves the problem. Namely, in this new model, the broken $\mathrm{U}(1)_Y$ (a subgroup of $\mathrm{SU}(5)_{GUT}$) is restored by quantum mechanical effects and hence we do not need to introduce the extra $\mathrm{U}(1)_H$. In fact, we find a stable quantum vacuum corresponding to the desired symmetry breaking $\mathrm{SU}(5)_{GUT} \rightarrow \mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$. We also show that there remains a pair of Higgs doublets massless in this quantum vacuum. In addition to these Higgs doublets, however, we have a pair of massless singlets which contributes to the renormalization-group flows of low-energy gauge coupling

constants and affects the unification of gauge couplings at the GUT scale. We also propose a simple solution to this problem, giving a sufficiently large mass to the singlet pair.

2. The model

Let us consider a supersymmetric hypercolor $\mathrm{SU}(3)_H$ gauge theory [4, 5] with 6 flavors of hyperquark chiral superfields Q_α^A and $\bar{Q}_A^\alpha$ $(\alpha = 1, \cdots, 3)$ in the fundamental representations $\mathbf{3}$ and $\mathbf{3}^*$ of $\mathrm{SU}(3)_H$, respectively. Here, A denotes flavor index and it runs from 1 to 6. The first five Q_α^I and $\bar{Q}_I^\alpha$ $(I = 1, \cdots, 5)$ transform as $\mathbf{5}^*$ and $\mathbf{5}$, respectively, under the GUT gauge group $\mathrm{SU}(5)_{GUT}$, while the last Q_α^6 and $\bar{Q}_6^\alpha$ are singlets of $\mathrm{SU}(5)_{GUT}$ [5].

We introduce three kinds of $\mathrm{SU}(3)_H$-singlet chiral superfields, a pair of H_I and $\bar{H}^I$, $\Sigma^I{}_J$ and Φ $(I, J = 1, \cdots, 5)$, which are $\mathbf{5}$, $\mathbf{5}^*$, $\mathbf{24}+\mathbf{1}$ and $\mathbf{1}$ of $\mathrm{SU}(5)_{GUT}$. Besides the gauge symmetry, $\mathrm{SU}(3)_H \times \mathrm{SU}(5)_{GUT}$, we impose a global $\mathrm{U}(1)_A$ symmetry [4, 5]

$$
\begin{array}{rcl}
Q_\alpha^I, \bar{Q}_I^\alpha & \rightarrow & Q_\alpha^I, \bar{Q}_I^\alpha, \\
Q_\alpha^6, \bar{Q}_6^\alpha & \rightarrow & e^{i\xi} Q_\alpha^6, e^{i\xi} \bar{Q}_6^\alpha, \\
H_I, \bar{H}^I & \rightarrow & e^{-i\xi} H_I, e^{-i\xi} \bar{H}^I, \\
\Sigma^I{}_J & \rightarrow & \Sigma^I{}_J, \\
\Phi & \rightarrow & e^{-2i\xi} \Phi, \\
(I = 1, \cdots, 5). & &
\end{array}
\tag{1}
$$

With the above gauge and global symmetries, we have a superpotential

$$
\begin{aligned}
W = \; & \lambda \Sigma^I{}_J \bar{Q}_I^\alpha Q_\alpha^J + h H_I \bar{Q}_6^\alpha Q_\alpha^I + h' \bar{H}^I \bar{Q}_I^\alpha Q_\alpha^6 + f \Phi \bar{Q}_6^\alpha Q_\alpha^6 \\
& + \frac{1}{2} m_\Sigma Tr(\Sigma^2) + \frac{1}{2} m'_\Sigma (Tr\Sigma)^2 - \mu_\Sigma Tr\Sigma \,.
\end{aligned}
\tag{2}
$$

Here, we have omitted trilinear self-coupling terms, $Tr(\Sigma^3)$, $Tr(\Sigma^2)Tr\Sigma$, and $(Tr\Sigma)^3$, for simplicity, and the mass term, $m_Q \bar{Q}_I^\alpha Q_\alpha^I$, has been absorbed into the Σ field by

shift of the trace part of Σ [1]. The terms, $\bar{H}^I H_I$, $\bar{Q}_6^\alpha Q_\alpha^6$ and $\Phi^n (n = 1, \cdots, 3)$, for example, are forbidden by the global $\mathrm{U}(1)_A$. Notice that the singlet Φ is not assumed in the previous models [4, 5]. As we will see later, Φ plays a crucial role on quantum mechanical restoration of the $\mathrm{U}(1)_Y$ gauge symmetry in the present model.

We first consider the following classical vacuum discussed in the previous analyses [4, 5]:

$$
\begin{aligned}
&\langle Q_\alpha^A \rangle = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ v & 0 & 0 \\ 0 & v & 0 \\ 0 & 0 & v \\ 0 & 0 & 0 \end{pmatrix}, \; \langle \bar{Q}_A^\alpha \rangle = \begin{pmatrix} 0 & 0 & v & 0 & 0 & 0 \\ 0 & 0 & 0 & v & 0 & 0 \\ 0 & 0 & 0 & 0 & v & 0 \end{pmatrix}, \\
&\langle \Sigma^I{}_J \rangle = \frac{\mu_\Sigma}{m_\Sigma + 2m'_\Sigma} \begin{pmatrix} 1 & & & & \\ & 1 & & & \\ & & 0 & & \\ & & & 0 & \\ & & & & 0 \end{pmatrix}, \\
&\langle H_I \rangle = \langle \bar{H}^I \rangle = 0 \,,
\end{aligned}
\tag{3}
$$

with $A = 1, \cdots, 6;\; I, J = 1, \cdots, 5$ and

$$
v = \sqrt{\frac{m_\Sigma \, \mu_\Sigma}{\lambda (m_\Sigma + 2m'_\Sigma)}} \,. \tag{4}
$$

Here, the vacuum-expectation value of Φ is undetermined since its potential is flat for $\langle Q_\alpha^6 \rangle = \langle \bar{Q}_6^\alpha \rangle = 0$. In this classical vacuum the gauge group is broken down as

$$
SU(3)_H \times SU(5)_{GUT} \to SU(3)_C \times SU(2)_L \,. \tag{5}
$$

There is no unbroken $\mathrm{U}(1)_Y$, which is the main motivation to introduce an extra $\mathrm{U}(1)_H$ gauge symmetry in Ref.[4, 5]. We now show that $\mathrm{U}(1)_Y$ broken in the classical vacuum is restored in a quantum vacuum.

[1]One may assume the hyperquark mass term, $m_Q \bar{Q}_I^\alpha Q_\alpha^I$, instead of introducing the trace part of Σ. For this case, one may also derive the same conclusion in the present paper.

3. Restoration of the $\mathrm{U}(1)_Y$ symmetry

Let us investigate the quantum vacua which satisfy the following conditions: (i) the vacuum-expectation values of Σ have a diagonal form with those of H and $\bar{H}$ vanishing, and (ii) at least two hyperquarks among Q^I_α and $\bar{Q}^\alpha_I$ $(I = 1, \cdots, 5)$ are massive in the vacua so that we can integrate out them to get a SUSY QCD-like theory with the effective $N_f = 4$ flavors. We identify the two massive hyperquarks with Q^I_α and $\bar{Q}^\alpha_I$ $(I = 1, 2)$. These quantum vacua include the vacuum defined in Eq.(3) in the classical limit.

We analyze the vacua in a similar way taken in Ref.[5]. The integration of massive hyperquarks, Q^I_α and $\bar{Q}^\alpha_I$ $(I = 1, 2)$, leads to a low-energy effective theory with a superpotential

$$\begin{aligned} W_{low} = \;& \lambda \Sigma^a{}_b \bar{Q}^\alpha_a Q^b_\alpha + h H_a \bar{Q}^\alpha_6 Q^a_\alpha + h' \bar{H}^a \bar{Q}^\alpha_a Q^6_\alpha \\ & + f \Phi\, \bar{Q}^\alpha_6 Q^6_\alpha - h h' \left(\frac{\bar{H}^1 H_1}{m_1} + \frac{\bar{H}^2 H_2}{m_2} \right) \bar{Q}^\alpha_6\, Q^6_\alpha \\ & + \frac{1}{2}\, m_\Sigma \Sigma^a{}_b \Sigma^b{}_a + \frac{1}{2}\, \frac{m_\Sigma m'_\Sigma}{m_\Sigma + 2m'_\Sigma} (\Sigma^a{}_a)^2 - \frac{m_\Sigma \mu_\Sigma}{m_\Sigma + 2m'_\Sigma} \Sigma^a{}_a \,, \end{aligned} \tag{6}$$

where $a, b = 3, 4, 5$, and m_1 and m_2 are the masses for the first and second hyperquarks, respectively [2]. Here, $\Sigma^I{}_J$ $(I, J = 1, 2)$ are decoupled from this low-energy superpotential and we have omitted $\Sigma^I{}_i$ and $\Sigma^i{}_I$ $(I = 1, 2,\ i = 3, \cdots, 6)$ since they are irrelevant for our analysis.

The effective superpotential which incorporates the strong $\mathrm{SU}(3)_H$ dynamics is described by meson $M^i{}_j$ and baryon B_i, $\bar{B}^i$ chiral superfields [7]

$$\begin{aligned} M^i{}_j &\sim Q^i_\alpha \bar{Q}^\alpha_j \,, \\ B_i &\sim \epsilon^{\alpha\beta\gamma} \epsilon_{ijkl} Q^j_\alpha Q^k_\beta Q^l_\gamma \,, \\ \bar{B}^i &\sim \epsilon_{\alpha\beta\gamma} \epsilon^{ijkl} \bar{Q}^\alpha_j \bar{Q}^\beta_k \bar{Q}^\gamma_l \,, \end{aligned} \tag{7}$$

[2] We have assumed $\langle Q^1_\alpha \bar{Q}^\alpha_1 \rangle = \langle Q^2_\alpha \bar{Q}^\alpha_2 \rangle = 0$, which is confirmed in view of the following analysis.

as follows:

$$
\begin{aligned}
W_{eff} = \quad & \Lambda^{-5}(B_i M^i{}_j \bar{B}^j - \det M^i{}_j) + \lambda \Sigma^a{}_b M^b{}_a \\
& + h H_a M^a{}_6 + h' \bar{H}^a M^6{}_a + f \Phi\, M^6{}_6 \\
& - h h' \left(\frac{\bar{H}^1 H_1}{m_1} + \frac{\bar{H}^2 H_2}{m_2} \right) M^6{}_6 \\
& + \frac{1}{2}\, m_\Sigma \Sigma^a{}_b \Sigma^b{}_a + \frac{1}{2}\, \frac{m_\Sigma m'_\Sigma}{m_\Sigma + 2m'_\Sigma} (\Sigma^a{}_a)^2 - \frac{m_\Sigma \mu_\Sigma}{m_\Sigma + 2m'_\Sigma} \Sigma^a{}_a \,,
\end{aligned}
\tag{8}
$$

where Λ denotes a dynamical scale of the low-energy $\mathrm{SU}(3)_H$ interactions, $a, b = 3, \cdots, 5$, and $i, j, k, l = 3, \cdots, 6$ [3].

For $\langle \Phi \rangle = 0$, we can see that B_6 and $\bar{B}^6$ necessarily have non-vanishing vacuum-expectation values leading to the breaking of $\mathrm{U}(1)_Y$ as in the classical vacuum in Eq.(3). However, there is a desired supersymmetric vacuum where

$$
\begin{aligned}
& \langle B_i \rangle = \langle \bar{B}^i \rangle = 0 \,, \\
& \langle M^6{}_6 \rangle = \langle M^6{}_a \rangle = \langle M^a{}_6 \rangle = 0 \,, \\
& \langle \Sigma^a{}_b \rangle = 0 \,, \\
& \langle M^a{}_b \rangle = \frac{m_\Sigma \mu_\Sigma}{\lambda (m_\Sigma + 2m'_\Sigma)}\, \delta^a{}_b \,, \quad \langle \Phi \rangle = \frac{1}{f \Lambda^5} \left[\frac{m_\Sigma \mu_\Sigma}{\lambda (m_\Sigma + 2m'_\Sigma)} \right]^3 \,, \\
& (i = 3, \cdots, 6 \text{ and } a, b = 3, \cdots, 5) \,.
\end{aligned}
\tag{9}
$$

Notice that in this vacuum the sixth hyperquark has an effective mass $m_6 = f \langle \Phi \rangle$ while other three hyperquarks are massless. Therefore, our model becomes a SUSY QCD-like theory with $N_f = 3$ massless quarks. The quantum vacuum in Eq.(9) is not the same as the classical one, which is consistent with the result obtained by Seiberg [7] for the case of $N_f = N_c$, with N_c being the number of colors.

With $\langle M^6{}_6 \rangle = 0$ we find that the Higgs doublets H_I and $\bar{H}^I$ $(I = 1, 2)$ remain massless in this quantum vacuum [5]. On the other hand, their color partners H_a and $\bar{H}^a$ have GUT-scale masses.

[3] The global $\mathrm{U}(1)_A$ has, of course, $\mathrm{SU}(3)_H$ gauge anomalies and it is broken by the nonperturbative effects. However, the effective superpotential W_{eff} in Eq.(8) except the dynamical part $(B_i M^i{}_j \bar{B}^j - \det M^i{}_j)$ respects the $\mathrm{U}(1)_A$ symmetry as noted in Ref.[5].

When $\langle \Phi \rangle \neq 0$ as in Eq.(9), the sixth hyperquark becomes massive. Therefore, the above vacuum in Eq.(9) may be also studied by means of the effective superpotential obtained through integrating out the sixth hyperquark instead of the second one:

$$\begin{aligned} W'_{eff} = \; & \Lambda'^{-5}(B_i M^i{}_j \bar{B}^j - \det M^i{}_j) + \lambda \Sigma^i{}_j M^j{}_i \\ & - \frac{hh'}{m_6} H_i \bar{H}^j M^i{}_j + \frac{1}{2} m_\Sigma \Sigma^i{}_j \Sigma^j{}_i + \frac{1}{2} \frac{m_\Sigma m'_\Sigma}{m_\Sigma + m'_\Sigma} (\Sigma^i{}_i)^2 - \frac{m_\Sigma \mu_\Sigma}{m_\Sigma + m'_\Sigma} \Sigma^i{}_i \,, \end{aligned} \tag{10}$$

where $i, j = 2, \cdots, 5$. The dynamical scale Λ' is different from the previous Λ by a simple-threshold relation [8] $m_2 \Lambda'^5 = m_6 \Lambda^5$. From this superpotential we obtain

$$\langle M^2{}_2 \rangle = 0, \quad \langle \Sigma^2{}_2 \rangle = \frac{\mu_\Sigma}{m_\Sigma + 2m'_\Sigma} \,. \tag{11}$$

Similarly we obtain $\langle M^1{}_1 \rangle = 0$ and $\langle \Sigma^1{}_1 \rangle = \langle \Sigma^2{}_2 \rangle$, which lead us to conclude that in the quantum vacuum in Eq.(9) we have the desired breaking of the GUT gauge group down to the standard-model one $SU(3)_C \times SU(2)_L \times U(1)_Y$.

As mentioned in the introduction, we have a pair of massless bound states B_6 and $\bar{B}^6$. Since they have $\mathrm{U}(1)_Y$-charges, they contribute to the renormalization-group equations of three gauge coupling constants in the standard model. We now show that a natural extension of our original superpotential given in Eq.(2) induces a sufficiently large mass to this dangerous singlet pair.

It seems quite reasonable that there are nonrenormalizable operators in the superpotential suppressed by some scale M_0 higher than the GUT scale (originating from gravitational interaction, for example). We, thus, consider the lowest-dimensional nonrenormalizable operator consistent with our gauge and global symmetries that is to contain baryon superfields. That is

$$\delta W = \frac{f'}{M_0^3} \epsilon^{\alpha\beta\gamma} \epsilon_{\alpha'\beta'\gamma'} (Q^I_\alpha Q^J_\beta Q^K_\gamma)(\bar{Q}^{\alpha'}_I \bar{Q}^{\beta'}_J \bar{Q}^{\gamma'}_K) \,. \tag{12}$$

This interaction generates a mass term for B_6 and $\bar{B}^6$ in the effective superpotential

as

$$\delta W_{eff} = \frac{f'}{M_0^3} B_6 \bar{B}^6 \,, \tag{13}$$

which corresponds to the physical mass for B_6 and $\bar{B}^6$

$$m_{B_6} \simeq \frac{f'\Lambda^4}{M_0^3} \,. \tag{14}$$

Taking $f' \sim \mathcal{O}(1)$, $\Lambda \sim 10^{16} - 10^{17}$ GeV and $M_0 \sim 10^{17} - 10^{18}$ GeV, we obtain $m_{B_6} \sim 10^{13} - 10^{14}$ GeV. In Fig.1 we show two-loop renormalization-group flows of three gauge coupling constants, α_1, α_2 and α_3, with $m_{B_6} = 10^{13}$ GeV and $\alpha_3(m_Z) = 0.117 \pm 0.010 (2\sigma)$ [9]. We see that our model is marginally consistent with the GUT unification and the small value of α_3 is favored. It may be interesting that the recently reported Zbb anomaly implies rather small $\alpha(m_Z)$ [10].

It is remarkable that the higher-dimensional operator in Eq.(12) stabilizes the present vacuum given in Eq.(9) and there is no flat direction around the vacuum besides the directions of $\langle H_I \rangle = \langle \bar{H}^I \rangle$ $(I = 1, 2)$ [4].

4. Conclusion

In this paper we have proposed a model for the dynamical generation of light Higgs doublets in SUSY-GUT's based on a semisimple gauge group SU(3)$_H \times$ SU(5)$_{GUT}$. The SU(3)$_H$ interactions become strong at the GUT scale and cause a dynamical breaking of the SU(5)$_{GUT}$. We have found that there is indeed a desired quantum vacuum corresponding to the breaking $SU(5)_{GUT} \to SU(3)_C \times SU(2)_L \times U(1)_Y$. We have also noted that an introduction of a higher-dimensional operator in Eq.(12) to the superpotential stabilizes the desired quantum vacuum. Moreover, the same operator gives rise to a mass for the composite baryon states B_6 and $\bar{B}^6$ of the order of $10^{13} - 10^{14}$ GeV, which makes our model consistent with the GUT unification of three gauge coupling constants of $SU(3)_C \times SU(2)_L \times U(1)_Y$.

[4]Without the new term in Eq.(12) we have another flat direction given by $\langle B_6 \rangle \langle \bar{B}^6 \rangle + f \langle \Phi \rangle = \det \langle M^a{}_b \rangle$.

Finally, we comment on the proton decay via the dimension-five operators [11]. In the present vacuum, the mass matrix for color-triplet states M, H and $\bar{H}$ is given by

$$(H_a \; M^6{}_a) \; \hat{m}_C \begin{pmatrix} \bar{H}^a \\ M^a{}_6 \end{pmatrix}, \tag{15}$$

where

$$\hat{m}_C = \begin{pmatrix} 0 & h \\ h' & \dfrac{v^4}{\Lambda^5} \end{pmatrix}. \tag{16}$$

with $v^2 = \dfrac{m_\Sigma \mu_\Sigma}{\lambda(m_\Sigma + 2m'_\Sigma)}$. The dimension-five operators are proportional to $\left(\hat{m}_C^{-1}\right)_{11}$, which is at the order of the GUT scale, $\sim \dfrac{v^4}{hh'\Lambda^5}$. Thus, we have unsuppressed proton decays [5]. This is a crucial point to distinguish the present model from the previous one [4, 5]. The reason for having non-vanishing dimension-five operators is that $\langle \Phi \rangle \neq 0$ induces the effective mass for Q^6_α and $\bar{Q}^\alpha_6$ and hence there remains no $\mathrm{U}(1)_A$ symmetry [6] , which forbids the dimension-five operators for the proton decays.

[5]From the present experimental limit on the proton decays we derive a constraint [12] $hh'\Lambda^5/v^4 \gtrsim 3 \times 10^{16}$ GeV.

[6]Notice that $\langle \Phi \rangle \neq 0$ never generates the mass term for H_I and $\bar{H}^I (I = 1, 2)$ as seen in Eq.(8).

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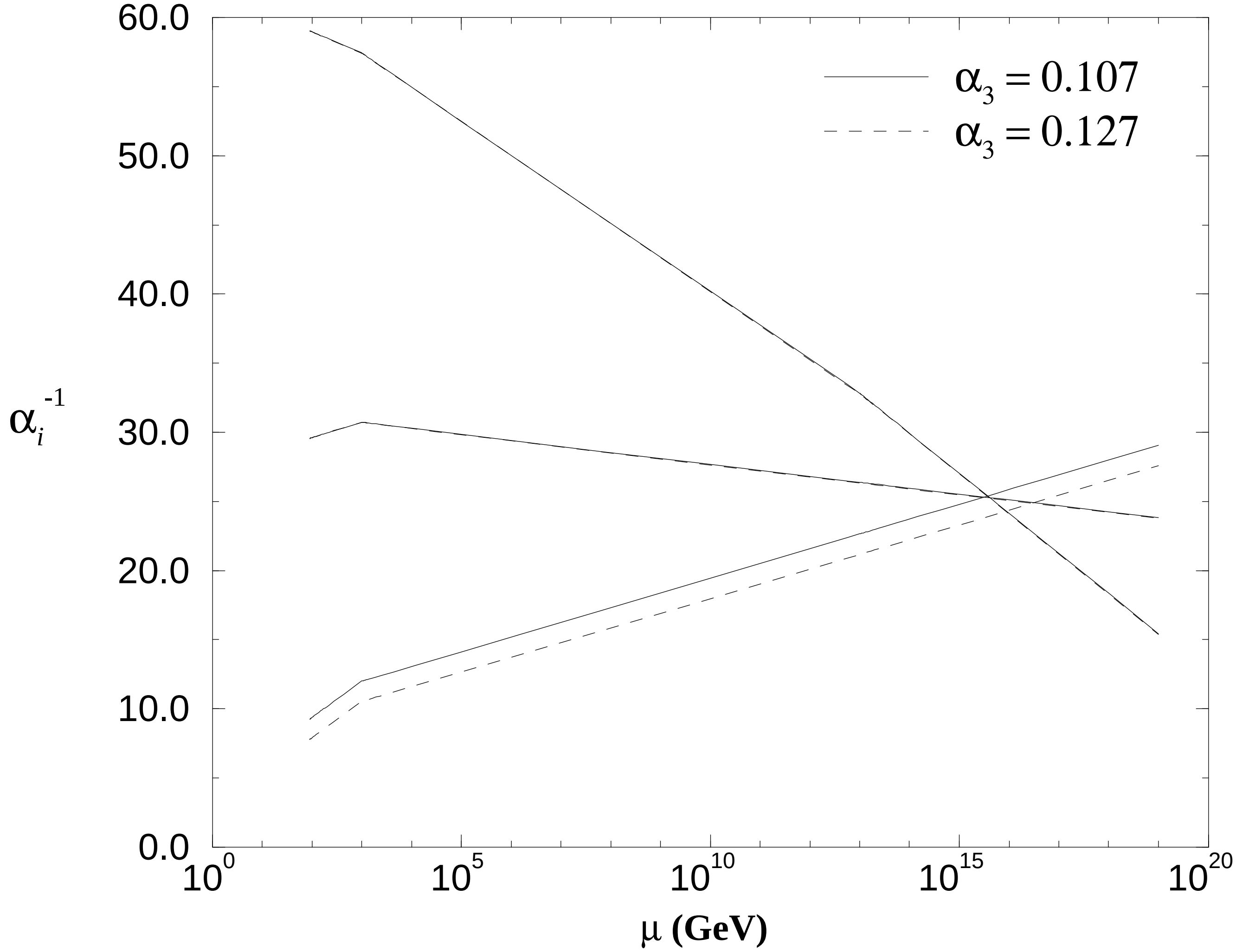
Figure caption

Fig.1 Two-loop renormalization-group flows of three gauge coupling constants, α_1, α_2 and α_3, for $m_{B_6} = 10^{13}$ GeV. The solid (dashed) lines are obtained for $\alpha_3(m_Z) = 0.107\,(0.127)$ which are within 2σ ranges from mean value [9]. We have assumed the SUSY-breaking scale to be 1 TeV.



60.0
50.0
40.0
30.0
20.0
10.0
0.0
α_i^{-1}
10^0
10^5
10^{10}
10^{15}
10^{20}
μ (GeV)
$\alpha_3 = 0.107$
$\alpha_3 = 0.127$